Multi-quark potential from AdS/QCD

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Abstract

Heavy multi-quark potential in the $SU(N)$ color group using hard-wall AdS/QCD at both zero and finite temperature is studied. A Cornell-like potential is obtained for baryons and other exotic configurations and compared with those in the quenched lattice calculation in $N = 3$ case. At the end we also discuss possible improvements in the UV region of potential.

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I. INTRODUCTION

When the mass of a heavy quark m_Q is much larger than the QCD scale Λ_{QCD}, the heavy $\bar{Q}Q$ bound states (quarkonium) are like hydrogen atoms where ordinary quantum mechanics can be applied. In the non-relativistic limit, it is possible to describe the interaction between Q and $\bar{Q}$ in terms of a local potential $E(L)$, where L is inter-quark separation. By solving the Schrödinger equation in three-dimensions, the quarkonium spectrum can be adequately characterized by $E(L)$ and m_Q. This potential approach to quarkonia has been applied to the study of Charmonium and Bottomonium states and a fair agreement with measurement was found. One of the most popular models for $E(L)$ is the so-called *Cornell* potential[1],

$$E(r) = -\frac{A}{L} + \sigma L. \tag{1}$$

The first part is Coulombic due to the one-gluon exchange between Q and $\bar{Q}$. The second part is the confinement potential due to the formation of color flux tube. Relativistic, radiative and other non-perturbative corrections can also be added *ad hoc* as extra terms. Though we are still lacking a first-principle derivation of the whole potential, among few approaches to the non-perturbative QCD, lattice QCD computation provides an irreplaceable way to test its validity.

On the other hand, the color flux tube can be easily modelled by a string. In particular, the coefficient σ in the equation (1) is identified as the string tension. Though this old version of String Theory first failed to be a realistic QCD model for its false prediction of some massless particles, it nevertheless took another ambitious task as a candidate of Theory of Everything and later evolved into its supersymmetric counterpart, the Superstring Theory.

However, not only QCD string survived but in fact it revived by gaining new insights from the so-called Anti-de Sitter space-Conformal Field Theory (AdS/CFT) correspondence, first proposed in [2, 3, 4]. In particular, it was applied to relate the thermodynamics of $\mathcal{N} = 4$ super Yang-Mills (SYM) theory in four dimensions to the thermodynamics of Schwarzschild black holes in five-dimensional Anti-de Sitter space[5]. In this description, confinement/deconfinement phase transition of gauge theory on a sphere has its holographic dual description as the Hawking-Page phase transition[6]. Later, it was realized that confinement can be achieved by capping off the Calabi-Yau cone smoothly at the infrared tip[7, 8, 9], or by introducing IR cutoff in the AdS space [10, 11, 12, 13, 14, 15, 16, 17]. This latter approach are usually referred to as the bottom-up construction of AdS/QCD,

together with the other top-down construction by [18, 19], are two main approaches to realize QCD physics via the (super)gravity theory in specific backgrounds. In particular, it is hoped that, as an alternative to lattice QCD, the non-perturbative region of QCD can be better understood both qualitatively and quantitatively within its dual picture of gravity. One important testing ground for the AdS/QCD is the experiments of Relativistic Heavy Ion Collision (RHIC), collisions of gold nuclei at 200 GeV per nucleon are about to produce a strongly-coupled quark-gluon plasma (QGP), which behaves like a nearly ideal fluid. While the perturbative calculation can not be fully trusted in this strongly-coupled region, there are increasing amount of interests in calculation of hydrodynamical transport quantities via the use of AdS/CFT correspondence, in particular that the energy loss of a heavy quark moving through $\mathcal{N} = 4$ SYM thermal plasma has been extensively studied[20]. The drag force was derived in the context of AdS/CFT to model the effective viscous interaction [21, 22], later it was generalized to a rotating black hole or with a dilaton field [23, 24, 25, 26] and B-field[27, 28]. Drag force of a comoving $(Q\bar{Q})$ pair was also considered in [29] and energy loss of baryon was studied in [30]. It is hoped that this line of research will eventually make contact with experimental results from RHIC.

Back to the equation (1), one would like to know how to derive it in the revived String Theory. We have learnt that $(Q\bar{Q})$ static potential can be calculated via temporal Wilson loop for time $T \gg L$,

$$< \mathcal{W} > \simeq e^{-TE(L)}. \tag{2}$$

In the AdS/QCD scenario, this quantity, at zero genus, can be calculated via on-shell classical (super)gravity action in the AdS bulk geometry, i.e. the minimal surface enclosed by the Wilson loop on the boundary[31]. Similarly, baryon potential can also be constructed once the baryon vertex is realized as a wrapped brane on the compactified sphere[32, 33, 34, 35, 36]. However, the potential obtained in the AdS background is always Coulombic thanks to its conformal symmetry. A Cornell-like potential for heavy $(Q\bar{Q})$ was obtained in [14] by breaking this symmetry via new scale set by an IR cut-off. This so-called hard-wall model simply puts an IR cut-off brane at location $r = R$. This cut-off brane is responsible for confinement as it will become clear later. The corresponding metric is then,

$$ds^2 = (\frac{r}{R})^2(-dt^2 + d\vec{x}^2) + (\frac{R}{r})^2 dr^2 + R^2 d\Omega_5^2 \tag{3}$$

where $R \leq r < \infty$.

In this paper, we go beyond the study of heavy $(Q\bar{Q})$ static potential, and look for multi-quark potential in the $SU(N)$ color group using the above-mentioned toy model of AdS/QCD at both zero and finite temperature.

This paper is organized as follows. In section II we refresh the construction of heavy $(Q\bar{Q})$ potential as in [14]. In section III we construct the heavy baryon potential and make a naive comparison with that in the quenched lattice calculation for $SU(3)$ color group[37]. In section IV we construct exotic multi-quark configuration. In particular, we obtain tetra-quark and penta-quark potential and compare them with the lattice results. In section V we study baryon potential at finite temperature. In section VI we make some proposals to improve our construction, but also discuss their limitations. In section VII, we conclude with summary and a few comments.

II. HEAVY MESON POTENTIAL

In this section, we recall the construction about inter-quark potential of heavy $(Q\bar{Q})$ in truncated AdS space. When the quark and anti-qaurk are close enough, the potential E and inter-quark distance L are given by

$$E(r_0) = \frac{r_0}{\pi\alpha'}\Big(\int_1^\infty dy\big(\frac{y^2}{\sqrt{y^4-1}} - 1\big) - 1\big),$$
$$L(r_0) = 2\frac{R^2}{r_0}\int_1^\infty dy\frac{1}{y^2\sqrt{y^4-1}}, \tag{4}$$

where r_0 is the lowest point where QCD string can reach in the bulk. The AdS (S^5) radius $R^4 = \lambda/\alpha'^2 = g_{YM}^2 N/\alpha'^2$. Therefore one obtains Coulomb-like potential[31]

$$E = -\alpha\frac{\sqrt{\lambda}}{L}, \qquad \alpha \simeq 0.228. \tag{5}$$

This result is expected from conformal invariance. As the distance increases, the U-shape string will be cut by the IR brane and the potential will include an additional term, given by[14]

$$E' = E(R) + \frac{1}{2\pi\alpha'}(L - L(R)). \tag{6}$$

As a result, we recover linear potential from the second term for large separation.
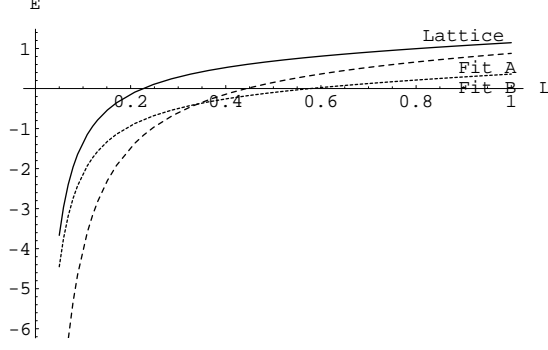
FIG. 1: Inter-quark potential as a function of baryon size. Fitting curve A agrees with the lattice simulation at IR region, while curve B agrees at UV region.

III. HEAVY BARYON POTENTIAL

The baryon $(Q \cdots Q)$ in the $SU(N)$ SYM is described by N open strings joining at a vertex formed by wrapping a D5 brane on S^5. The introduction of open strings is essential to incorporate quarks in fundamental representation of the SYM. The total energy of this baryon configuration is simply a sum of energy of N F1's and a wrapped D5 brane, given by[33]

$$E = -\beta N \frac{\sqrt{2\lambda}}{L}, \qquad \beta \simeq 0.007 \tag{7}$$

where the vertex is situated at the center of AdS_5 and the total energy has been regularized by subtracting off the energy of free quarks.

With the IR cut-off, we simply allow the configuration of each open string to be the same as in the original AdS space as if the vertex were at $r = r_0 < R$. However, those strings are cut by the IR brane and replaced by string segments on the brane. Therefore the vertex is actually projected onto the brane. The relation between r_0 and baryon radius L is still valid, i.e.

$$L = \frac{R^2}{r_0} \int_1^\infty \frac{dy}{y^2 \sqrt{\beta^2 y^4 - 1}} = \gamma \frac{R^2}{r_0}, \qquad \gamma \simeq 0.481 \tag{8}$$

and the projected radius L_R on the cut-off is

$$L_R = \frac{R^2}{r_0} \int_1^{y_R} \frac{dy}{y^2 \sqrt{\beta^2 y^4 - 1}}, \tag{9}$$

where $y \equiv \frac{r}{r_0}$, $y_R \equiv \frac{R}{r_0}$ and $\beta = \sqrt{16/15}$. The total energy, composed of one vertex and N
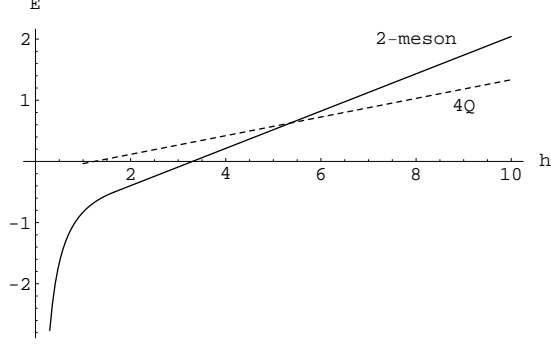
FIG. 2: Flux-tube recombination between connected four-quark state and two-meson state is referred to the *flip-flop*. It happens as two vertices are close enough.

strings, becomes

$$E = \frac{N}{2\pi\alpha'}r_0\Big(\int_{y_R}^{\infty} dy(\frac{\beta y^2}{\sqrt{\beta^2 y^4 - 1}} - 1) - 1\Big) + \frac{N}{2\pi\alpha'}L_R + \frac{R^5}{8\alpha'}. \qquad (10)$$

After replacing r_0 and E in terms of L, and choosing y_R properly, it becomes

$$E = -N\frac{A}{L} + \sigma NL + C. \qquad (11)$$

In particular, at the limit $r_0 \to 0$, one obtains[1]

$$A = \gamma\frac{\sqrt{\lambda}}{2\pi}, \qquad \sigma = \frac{1}{2\pi\alpha'}, \qquad C = \frac{R^5}{8\alpha'} \qquad (12)$$

In the FIG. 1, we compare with the quenched lattice result for $SU(3)$ color group[37], i.e. $(A, \sigma, C) = (0.0768, 0.1524, 0.9182)$. However, two variables are generally insufficient to fit three unknowns. If we fit the same σ and C with the choice of $\alpha' = 1.0443$ and $R^2 = 2.2591$ (or $\lambda = 5.1036$), then a twice bigger $A = 0.1656$ is obtained for curve A. If we fit the same A and σ, then we obtain $C = 0.1345$ (or $\lambda = 1.0064$) for curve B. As a result, fitting curve A agrees with the lattice simulation in the IR region, while curve B agrees in the UV region.

IV.   EXOTIC MULTI-QUARK POTENTIAL

One can imagine that multi-quark system might be formed in the process of pair annihilation as two or more baryons/anti-baryons are dragged toward each other. In the case

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[1] L here refers to the projected distance from constituent quark to the vertex, not the inter-distance between quarks as used in the lattice simulation[37]. If the convention of lattice is used, we obtain $A' \simeq 0.133\sqrt{\lambda}$, about half the coefficient of meson.
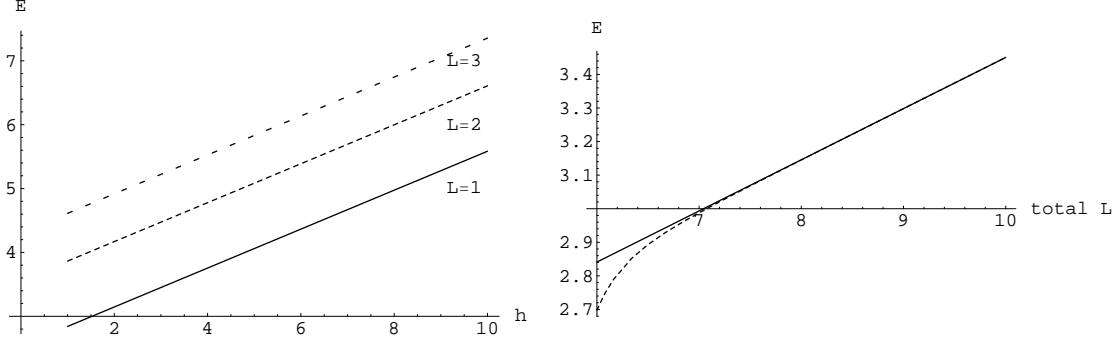
FIG. 3: To the left: It shows that penta-quark potential is a linear function of vertex separation h for various quark-vertex distance L. To the right: It shows that penta-quark potential is a linear function of total flux tube length(solid line). The dashed line indicates a deviation from linearity in the UV region thanks to either the correction from attraction between vertices or the flip-flop mechanism, as also predicted in the lattice simulation.

of $N = 3$, for instance, a baryon made of three quarks (QQQ) can combine with an anti-baryon made of three anti-quarks $(\bar{Q}\bar{Q}\bar{Q})$ to form a tetra-quark $(QQ\bar{Q}\bar{Q})$, by annihilating a pair of $(Q\bar{Q})$ and reconnecting two baryon vertices. Similarly, a penta-quark $(QQ\bar{Q}QQ)$ can be constructed by an anti-baryon and two baryons. In general, one could form a bound state of $(mN - 2m + 2)$ (anti-)quarks with m vertices as long as its energy is lower than sum of m (anti-)baryons[2]. From now on, we restrict ourselves to the case of $N = 3$. It is straightforward to write down inter-quark potential of a tetra-quark, i.e.

$$E_{4Q} = -4\frac{A}{L} + (4L + h)\sigma + 2C, \tag{13}$$

where h is the separation between two vertices. For small h, however, the lattice data tend to agree with the potential of two-meson, i.e. twice the potential given by (6) with separation h. This flux-tube recombination between connected four-quark state and two-meson state is referred to the *flip-flop*. The FIG. 2 shows that this flip-flop can also be realized in our model. Next, inter-quark potential of a penta-quark reads,

$$E_{5Q} = -5\frac{A}{L} + (4L + 2h)\sigma + 3C, \tag{14}$$

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[2] Here we simply mention those bound states made of single open chain of baryon vertices. In addition, there might be states with closed chain(s). We thank Kazuyuki Furuuchi to point out this.
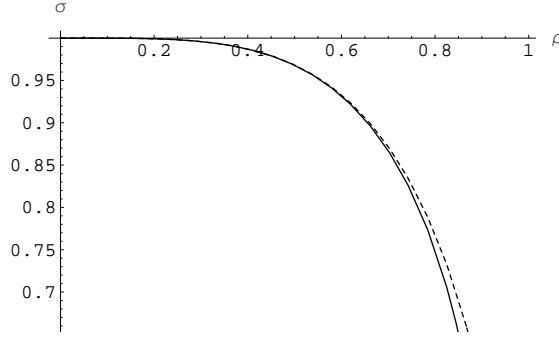
FIG. 4: The solid curve shows the temperature dependence of string tension (normalized to unity). It can be very well fit by the dashed curve given by equation (18) for low temperature.

where we have assumed that three vertices are placed in equal separation h. From the equation above, it is obvious to expect linear relation between potential E and separation h as shown in the left plot of FIG. 3, or between E and total length of color flux tube, say $4L + 2h$, as shown in the right plot of FIG. 3.

V. POTENTIAL AT FINITE TEMPERATURE

Now we are ready to study the same potential at the finite temperature, which corresponds to introducing an AdS black hole, i.e.

$$ds^2 = (\frac{r}{R})^2(-f(r)dt^2 + d\vec{x}^2) + (\frac{R}{r})^2\frac{1}{f(r)}dr^2 + R^2d\Omega_5^2, \tag{15}$$

where $f(r) = 1 - r_T^4/r^4$ and the event horizon is related to the Hawking temperature by $r_T = \pi R^2 T$. In the deconfined phase where $r_T \geq R$, the vertex falls into the black hole and leaves behind N dissociated quarks with total mass[3],

$$E = -N\frac{r_T}{2\pi\alpha'}. \tag{16}$$

However, the observers living in the boundary still see the vertex since it takes infinite time for the vertex to cross the horizon. In the confining phase where $r_T < R$, one may either argue that black hole is unaccessible to us so the potential has no temperature dependence

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[3] This bare mass is not physical but by convention of renormalization since ideal heavy quarks should have infinite mass which is unmeasurable.

for thermal AdS, or one may pretend that black hole still curves the bulk geometry. In the latter situation, the baryon potential has the following linear component[33],

$$N\frac{\sqrt{\lambda}}{2\pi r_0}\int_1^{y_R}dy\sqrt{\frac{15-18\rho^4-\rho^8}{(y^4-\rho^4)(16y^4-15+2\rho^4+\rho^8)}},\tag{17}$$

where $\rho = r_T/r_0$. Here, different from the situation at zero temperature, it is not allowed to send $r_0 \to 0$. In fact, the reality condition requires that $r_0 > r_T$. In the FIG. 4 we see the temperature dependence of the effective string tension, which can be best fit in low temperature by

$$\frac{\sigma(T)}{\sigma(0)} \simeq \sqrt{1-\rho^4}.\tag{18}$$

VI. IMPROVEMENTS

We have argued in section III that our model can not completely fit the lattice result. In this section, we provide three possible improvements to the potential but also discuss their limitations.

A. UV perfection

As mentioned before that our model cannot completely fit the lattice result, introducing another UV cut-off r_c may resolve this. However, this implies that instead of our original assumption of (infinitely) heavy quarks, the Cornell potential is realized via quarks with a finite mass

$$m_Q \simeq \frac{r_c - R}{2\pi\alpha'}.\tag{19}$$

If this is the case, we should also consider string breaking (pair production) as inter-quark separation grows. A careful examination excludes the possibility of sending $r_0 \to 0$, which leaves more than one choice for combination of r_0 and r_c, or equivalently y_R and $y_c = r_c/r_0$. In fact, we only find reasonable solutions as $y_R < 2$. One of these choices is that, for instant, $y_R \sim 1.50$ and $y_c \sim 3.29$.

B. IR refinement

We start with a truncated model with IR cut-off at $r = R$. This implies that the confinement/deconfinement transition temperature is at $T_c = 1/\pi R$, which is $\simeq 212$ MeV for the first fit A and $\simeq 312$ MeV for the second fit B. However, this model may be refined with a different IR cut-off r_{IR}. This provides a new parameter for perfect fit, where

$$C = \frac{R^4}{\alpha'} r_{IR}. \tag{20}$$

with the choice of $r_{IR} \simeq 6.9891$, However, this cut-off corresponds to the transition temperature[15],

$$T_c = \frac{2^{1/4}}{\pi r_{IR}} \simeq 54 \quad MeV, \tag{21}$$

which is, unfortunately, far too low in comparison with the lattice prediction $T_c \simeq 190$ MeV[38].

C. Correction via closed string channel and flip-flop mechanism

The lattice simulation indicates that in the FIG. 3, relation between potential and length is not completely linear, especially for small separation h. The residual force outside each interacting group of quarks may help explaining the deviation from linearity. This could be realized in our model by two means: one way is to include the graviton (closed string) exchange between vertices. The other is the flip-flop mechanism in short h. Both give the expected $1/h$ behaviour in the IR region.

VII. CONCLUSION

In this paper, we have constructed static potential of heavy baryon and other exotic multi-quark configurations in the hard-wall model of AdS/QCD. In particular, we obtained a Cornell-like potential and make naive comparison with that in the quenched lattice calculation for $SU(3)$ color group[37]. We found that the flux tube recombination between four-quark and two-meson states (flip-flop) happens when two vertex approaches each other in a tetra-quark configuration. In a penta-quark configuration, static potential is always linear with respect to inter-vertex distance or total length of flux tube, though a deviation

from linearity is expected for short h. We later proposed three different ways to improve our model. For UV perfection, we introduced a UV cut-off and there were more than one choices to achieve the goal. Nevertheless, the quark mass became finite and one has to take into account pair production (string breaking) as inter-quark distance grows. For IR refinement, we traded a tunable IR cut-off for perfect fit. However, this modification gave us far too low confinement/deconfinement transition temperature and was not so impressive. Finally, we argued that the deviation from linearity in the UV region in the figure (III) may come from correction via vertex interaction or flip-flop mechanism when the potential of two-meson state dominates. At end end, we would like to make a few remarks. First of all, there is no obvious reason to believe that AdS/QCD construction is trusted for large N as well as in the AdS/CFT correspondence. On the top of that, we should not ask too much for quantitative accuracy even for small N, especially without taking in account the $1/N$ correction. This may well explain that we did not find a perfect fit without causing another drawback. Secondly, in the confined phase, the string tension is expected to have a temperature correction $\propto -T^2$ instead of equation (18), observed by [14] as well. It is interesting to see if we could settle this disagreement in the context of AdS/QCD. Thirdly, there was early debate on Y-shape or Δ-shape flux tube configuration in the baryon ground state. Our construction always prefers the former since the latter costs more energy for slightly longer tube length and two extra vertex. At last, our proposal of UV perfection has similar effect when flavored quark is introduced[4], and a warping geometry with linear dilaton is responsible to the string breaking[39]. It would be interesting to realize this breaking scenario in the soft-wall QCD model. In conclusion, the new feature as well as advantage of this construction, in comparison with the old version QCD string, is to realize all parts of Cornell potential, though not perfect, by one piece of configuration of strings and vertices stretching over truncated AdS bulk. This implies that there may be a chance in AdS/QCD construction for a unified theory of QCD in both perturbative and non-perturbative regions.

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[4] The author would like to thank Carlos Nunez for pointing out this similarity.

Acknowledgments

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